

Ineffective Supersymmetry: Electroweak Symmetry Breaking from Extra Dimensions

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Abstract

Recently, a mechanism for electroweak symmetry breaking (EWSB) was discussed [1], in which the scale of EWSB is set by the scale of an additional dimension $R \sim \text{TeV}^{-1}$. The mechanism involves supersymmetry, but broken in such a fashion that high (four-dimensional) momentum loops are cut off by the finite size of the radius. In a Kaluza-Klein decomposition, a hard cutoff seems to give a strong cutoff dependence, while summing the entire tower is not only cutoff insensitive, but actually finite. Such behavior is easily understood in a formulation that respects five-dimensional locality. Finally, we note that certain models of this type naturally give operators which can “fake” the presence of a light Higgs in precision electroweak observables.

1 Introduction

A central question in particle physics is this: what is the origin of the scale of electroweak symmetry breaking, and what stabilizes it against radiative corrections? In the past two decades, a tremendous amount of effort has gone into answering this question, and exploring the implications of possible answers. This “hierarchy problem” is summed up simply: why is $M_W \ll M_{Pl}$?

Numerous proposals have been made to stabilize the weak scale against radiative corrections, two notable examples being technicolor [2] and weak-scale supersymmetry [3]. In technicolor, divergences associated with the Higgs sector are controlled because the Higgs is not a fundamental field in the theory. In supersymmetric theories, above the supersymmetry breaking scale, bosonic and fermionic loops cancel, leading to relative insensitivity to the cutoff of the theory.

More recently, the question of the hierarchy problem has been recast in theories with large extra dimensions [4, 5]. It was noted that if there are n extra dimensions, the effective, four-dimensional Planck scale is related to the higher dimensional Planck scale M_* , by the equation

$$M_{Pl}^2 = M_*^{2+n} V, \tag{1}$$

where V is the volume of the additional n dimensions. If this volume is large, M_* can be brought down nearly to the weak scale. If this is the case, loop divergences of the standard model are cut off at the TeV scale. While this is not a solution to the hierarchy problem, it recasts it as the question of the origin of the large volume of the extra dimensions.

While these proposals offer interesting frameworks within which to work, none of them is complete without an understanding of the origin of electroweak symmetry breaking. For instance, supersymmetry with gravity mediated supersymmetry breaking [6] offers an attractive understanding of symmetry breaking. If we assume general soft breaking masses defined at $\mu = M_{Pl}$, top/stop loops drive the Higgs soft mass-squared negative in the RG evolution, triggering EWSB at roughly the scale of supersymmetry breaking.

In contrast, there is no obvious explanation of EWSB in theories with large extra dimensions. Various limits require the higher dimensional Planck scale to be relatively large compared to M_W , with $M_* \sim 10\text{TeV}$. Without a theory of electroweak symmetry breaking, it is impossible even to begin to discuss whether this scale is unnaturally high, or, equivalently, to understand when limits on M_* become meaningful.

In this paper, we will employ a new mechanism for electroweak symmetry breaking appropriate for large extra dimension (LED) theories. It will require the presence of an additional dimension of size $R^{-1} \sim \text{TeV}$, and supersymmetry in the full five-dimensional theory. It will be different from previous

models with TeV sized extra dimensions [7, 8, 9, 10, 11, 12, 13], in that both gauge fields and matter fields - Q, U, D, L and E - will all propagate in five dimensions. Yukawas will be present on one brane, the “Yukawa brane”, while supersymmetry breaking will arise either localized at some distance from the Yukawa brane (the SUSY-breaking brane), or through global properties of the five dimensional space, as in the Scherk-Schwarz mechanism [14]. Because high momentum loops will be unable to sense both supersymmetry breaking and the Yukawas simultaneously, contributions to the Higgs mass will be insensitive to the cutoff of the theory. In fact, in contrast with ordinary four-dimensional supersymmetry theories, in these theories the Higgs soft mass will be *finite* in the UV!

1.1 Supersymmetry and Extra Dimensions?

At first glance, it may appear that using both supersymmetry and large extra dimensions to address the hierarchy problem is somehow redundant. After all, isn’t supersymmetry on its own a solution to the hierarchy problem?

The answer, of course, is that *weakly broken* supersymmetry is a solution to the hierarchy problem, rather than supersymmetry itself. Given a supersymmetric theory, one needs an additional sector to generate an exponentially small scale for supersymmetry breaking. Only then does one have a complete solution to the hierarchy problem.

Here, we would argue that we have merely replaced sector generating the small scale of supersymmetry breaking with a sector which generates an exponentially large volume for the extra dimensions, or, equivalently, the exponentially small scale $V^{-1/n}$. Then supersymmetry can be broken at order one.

Moreover, part of our motivation for considering extra dimensions is the presence in superstring theory. Since supersymmetry appears in these theories, it is worthwhile to consider the effects of supersymmetry in LED theories.

2 Features of Five-Dimensional Theories

Let us briefly review the standard formalism for approaching five dimensional theories.

To begin with, let us consider the Klein-Gordon equation for a scalar field in five dimensions:

$$(\partial_t^2 - \partial_{\vec{x}}^2 - \partial_y^2)\phi(t, \vec{x}, y) = 0, \quad (2)$$

where y is the coordinate of a compact fifth dimension, with radius R . Because the fifth dimension is compact, the values for the momentum in the fifth direction can only take on discrete values,

$0, R^{-1}, 2R^{-1}$, etc. When we Fourier decompose $\phi(t, \vec{x}, y)$, the Klein-Gordon equation becomes

$$(\partial_t^2 - \partial_{\vec{x}}^2 + k^2 R^{-2})\phi_k(t, \vec{x}) = 0. \quad (3)$$

Thus each Fourier mode (or, equivalently, each Kaluza-Klein mode), acts as a four-dimensional field with mass k/R . The y dependence has been replaced by the values of an infinite tower of fields. Although we have used scalars here as an example, the same is true for higher spin particles.

If additional dimensions are present, we expect to see Kaluza-Klein towers for the fields which propagate in them. For standard-model fields, the absence of four-fermion operators, and corrections to precision electroweak observables generally constrain these dimensions to be somewhat small $R^{-1} \sim \text{TeV}$ [15]. In contrast, gravity can propagate in dimensions as large as $R \sim 10^{-5}m$ [16].

2.1 Supersymmetry in Five Dimensions

If we wish to discuss five dimensional supersymmetric theories, we need to understand the differences between four- and five-dimensional supersymmetry. Supersymmetry has a fermionic generator, and the minimal fermionic representation of the five-dimensional Lorentz group is a four-component spinor. This spinor decomposes under the four-dimensional Lorentz group as two two-component Weyl spinors. Thus, we see that five-dimensional $N = 1$ supersymmetry is equivalent to four-dimensional $N = 2$ supersymmetry.

In $N = 2$ supersymmetry, the minimal representations of the supersymmetry algebra have more components than in $N = 1$ [17]. For instance, chiral superfields become part of hypermultiplets which contain two chiral superfields. Moreover, from a four-dimensional point of view, these superfields are vectorlike under gauge symmetries. That is, given a hypermultiplet (Φ, Φ^c) , if Φ transforms under $SU(n)$ as an \mathbf{N} , then Φ^c transforms as an $\overline{\mathbf{N}}$.

Similarly, gauge field multiplets no longer consist merely of a vector and a fermion. They now contain an additional chiral superfield which transforms as an adjoint under the gauge group.

In addition to enlarging the field content of the theory, $N = 2$ supersymmetry strongly constrains the form of the interactions we can write. In particular, we cannot write trilinear couplings, such as Yukawas. That is, terms like

$$\int d^4x dy d^2\theta f_t Q U H \quad (4)$$

are forbidden.

All of these issues seem problematic phenomenologically. For one thing, we know that matter is chiral, not vectorlike under the gauge groups. Where are these Q^c, U^c, D^c, L^c and E^c states? We

do not see massless chiral adjoint fields, and the masses of fermions seem to require the presence of trilinear interactions. How can we resolve these problems?

Let us separate these into separate issues: missing states and needed couplings. The missing states are typically removed by use of an “orbifold projection”. Simply put it is this: our Kaluza-Klein (Fourier) decomposition can be written in terms of even functions (cosines) and odd functions (sines). It is a consistent projection to require that both the conjugate states and chiral adjoint fields satisfy odd boundary conditions (i.e., only keep the sines), while the unconjugated fields and vector supersuperfields satisfy even boundary conditions (i.e., only keep the cosines). Then there are no massless chiral adjoints, and there are no massless partners of the chiral matter fields.

Under this orbifold projection, two points on the circle are special, $y = 0$ and $y = \pi R$. These “fixed points” are mapped to themselves under the transformation $y \rightarrow -y$. We can naturally place branes at these points, and the interactions on these branes need only preserve four-dimensional Lorentz invariance, and hence, only $N = 1$ supersymmetry. That is, while eq. 4 is forbidden, a term

$$\int d^4x dy d^2\theta \delta(y) f_t Q U H \quad (5)$$

is allowed.

3 Breaking Supersymmetry

We can now explore the interesting possibility: what if supersymmetry is broken “away” from the Yukawa brane? This can take different forms, and let us consider two distinct possibilities. First, perhaps supersymmetry is broken by an F-component expectation value of some chiral superfield X localized on the brane at $y = \pi R$. The second possibility is that we use the Scherk-Schwarz mechanism, and require different boundary conditions for particles and their superpartners.

We will focus our attention on the localized supersymmetry breaking scenario, and return to the case of Scherk-Schwarz models shortly.

Suppose that $F_X \sim M_*^2$ is localized at $y = \pi R$. Then interactions (setting $M_* = 1$)

$$\int d^4x dy d^4\theta c_X X^\dagger X Q^\dagger Q \delta(y - \pi R) \quad (6)$$

(and similar terms for U, D, L, E) will generate soft masses for superpartners. How large are these masses? If we use the naive Kaluza Klein decomposition, we get mass terms

$$\int d^4x c_X F_X^2 R^{-1} \tilde{q}_i^* \tilde{q}_j. \quad (7)$$

Focusing on the diagonal ($i = j$) entries of this, we might expect that all superpartners have masses $\sim F_X \sqrt{c_X R^{-1}}$ (again in units with $M_* = 1$). However, this assumes that such an operator is a small perturbation on the system and it is not. To find the mass spectrum, we should instead solve the differential equation

$$-\partial_y^2 g(y) + m^2 g(y) + c_X F_X^2 \delta(y - \pi R) g(y) = 0. \quad (8)$$

Solving this for even $g(y)$, we have solutions

$$g(y) = C \cos(my) \quad (9)$$

where m is a solution of the equation

$$\tan(m\pi R) = \frac{c_X F_X^2}{2m}. \quad (10)$$

For $R \gg 1$ in fundamental units, this has solutions

$$m \simeq (n + \frac{1}{2}) R^{-1}. \quad (11)$$

That is, we have shifted the Kaluza-Klein spectrum of our superpartners up one-half unit in R^{-1} ¹. In fact, at this (leading) order in R^{-1} , the spectrum is independent of the size c_X so long as it is order one! Flavor changing effects arising from superpartner loops are naturally suppressed by this high degree of degeneracy. In fact, it is easy to understand this behavior: the large mass term in eq. 6 attempts to give large masses to the superpartners. By avoiding the brane at $y = \pi R$ (that is, by developing a node at $y = \pi R$), the superpartners pick up masses $m \sim R^{-1}$ from the kinetic terms, but no mass from the brane. So long as $R^{-1} < F_X \sqrt{c_X R^{-1}}$, or, equivalently, $R^{-1} < F_X^2 c_x$, the masses of the lightest modes are smaller with an approximate node at the brane at $y = \pi R$.

Now that we have found the spectrum of our superpartners, we can calculate loops corrections to the soft mass of the Higgs field, which we show in figure 1. Note the presence of the new diagram involving the conjugate fields. This diagram is necessary to ensure a cancellation in the supersymmetric limit (and is easily derived from the boundary Yukawa with a bulk superpotential term $\int dy Q^c \partial_y Q$ [17, 18]).

We now have the spectrum of the theory, but to calculate the loop diagrams, we must properly normalize the modes. The wavefunctions $g_k(y)$ for the KK modes are normalized such that $\int_0^{\pi R} [g_k(y)]^2 dy = \pi R$. If η_k^i is the value of the k -th Kaluza Klein mode of the field i on the Yukawa

¹The analysis for gauginos is somewhat different, and is carried out in [1].

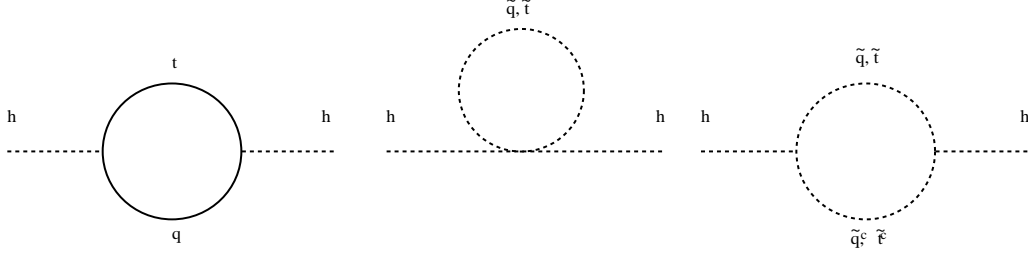


Figure 1: One-loop diagrams contributing to the Higgs-boson mass.

brane, we have

$$\eta_k^\psi = \left(\frac{1}{\sqrt{2}}\right)^{\delta_{k,0}}, \quad (12)$$

$$\eta_k^\phi = 1, \quad (13)$$

$$\eta_k^F = \begin{cases} \left(\frac{1}{\sqrt{2}}\right)^{\delta_{k,0}} & \text{for } r^F = 0, \\ 1 & \text{for } r^F = \frac{1}{2}, \end{cases} \quad (14)$$

where $k = 0, 1, 2$, etc.

With these normalizations, we can calculate the loop contributions to the Higgs soft mass. Using the variable $x = p_E R$, one can show

$$\begin{aligned} -i m_{\phi_H}^2 &= \frac{i N_c f_t^2 \epsilon^2}{R^2} \int \frac{d^4 x}{(2\pi)^4} x^2 \\ &\times \sum_{k,l=0}^{\infty} \left[\frac{(\eta_k^\psi)^2 (\eta_l^\psi)^2}{(x^2 + k^2)(x^2 + l^2)} - \frac{(\eta_k^\phi)^2 (\eta_l^F)^2}{(x^2 + (k + r^\phi)^2)(x^2 + (l + r^F)^2)} \right]. \end{aligned} \quad (15)$$

One might worry that each mode would contribute to the Higgs soft mass, and thus give a severe cutoff dependence. However, this is not the case: we can in fact sum the entire tower and achieve the *finite* result

$$m_{H_u}^2 = -\frac{3\zeta(3)}{8\pi^4} \frac{N_c y_t^2}{R^2}. \quad (16)$$

This result seems remarkable! It is especially so when we note that the lightest stops do not have the couplings appropriate for an ordinary four-dimensional SUSY theory. That is, if one only saw the lightest sfermions, their couplings to the Higgs differ from those expected from four-dimensional $N = 1$ supersymmetry by a factor $\sqrt{2}$. It is tempting to refer to this scenario as “ineffective supersymmetry.” While supersymmetry is protecting the Higgs mass, at no energy do we ever encounter a four-dimensional effective supersymmetric theory.

Here we have summed the whole tower. If we had employed a hard cutoff in 15, we would have found a severe cutoff dependence, in accord with naive expectations. How can we be certain that this technique (summing to infinity) is the right one?

Within the Kaluza-Klein formalism, it is somewhat difficult to see that this is correct. In fact, there has been a great deal of discussion recently [19, 20, 21, 22], as to whether such a summation amounts to a hidden fine tuning. What can we say?

A hard cutoff is clearly inappropriate. Since the physical features of the theory should determine its UV properties, any regulator we use must preserve certain features, namely, supersymmetry, five-dimensional Lorentz invariance, and, importantly, locality. To understand the finite nature of the result, we should formulate the calculation in which locality is manifest. Indeed, in the next section, we shall see from a five-dimensional perspective that the result *must* be finite.

4 Five dimensional interpretation

It is easiest to see that the result must be finite in a formalism which makes the fifth dimension explicit [1]. Thus, we shall work in mixed position-momentum space [23]. Since contributions to the Higgs soft mass rely both on Yukawas localized at $y = 0$ and supersymmetry breaking at $y = \pi R$, high momentum loops cannot simultaneously “see” both elements, causing an exponential damping at high pE .

One can calculate the scalar, fermion and F-component propagators easily [1]. In the large susy breaking limit, they are

$$\tilde{G}_\phi(k_4, y) = \frac{1}{2k_4} \frac{1}{\sinh[k_4\pi R]} \left\{ \cosh[k_4(\pi R - y)] - \frac{m \cosh[k_4 y]}{2k_4 \sinh[k_4\pi R] + m \cosh[k_4\pi R]} \right\}, \quad (17)$$

for $(y \in [0, \pi R])$. Again, we can take the large susy breaking limit, $m \rightarrow \infty$ and then the propagator becomes

$$\tilde{G}_\phi(k_4, y) = \frac{1}{2k_4} \frac{1}{\sinh[k_4\pi R]} \left\{ \cosh[k_4(\pi R - y)] - \frac{\cosh[k_4 y]}{\cosh[k_4\pi R]} \right\}, \quad (18)$$

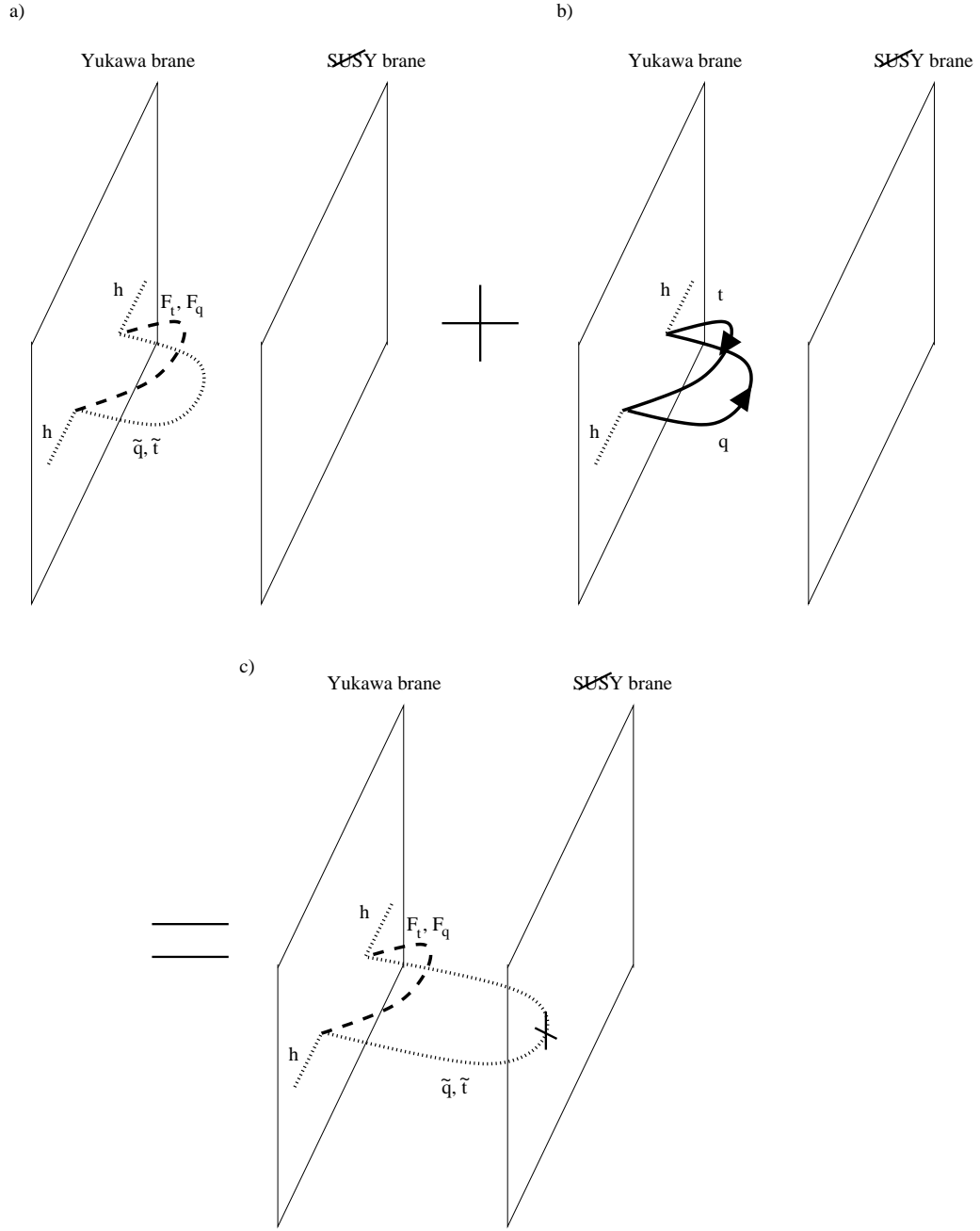
$$\tilde{G}_\psi(k_4, y) = \frac{k_4}{2k_4} \frac{\cosh[k_4(\pi R - y)]}{\sinh[k_4\pi R]}, \quad (19)$$

and

$$\tilde{G}_F(k_4, y) = \frac{k_4}{2} e^{-k_4|y|}. \quad (20)$$

The bosonic amplitude from Fig. 2(a) is given by

$$m_{\text{boson}}^2 = 2N_c \int \frac{d^4 k_4}{(2\pi)^4} (2\pi R y_t)^2 \frac{\tanh[k_4\pi R]}{2k_4} \frac{k_4 \coth[k_4\pi R]}{2}, \quad (21)$$



and the fermionic one from Fig. 2(b) by

$$m_{\text{fermion}}^2 = -N_c \int \frac{d^4 k_4}{(2\pi)^4} (2\pi R y_t)^2 \text{Tr} \left[\frac{k_4 \coth[k_4 \pi R]}{2k_4} \frac{(1 - \gamma_5)}{2} \frac{k_4 \coth[k_4 \pi R]}{2k_4} \frac{(1 + \gamma_5)}{2} \right]. \quad (22)$$

These are combined to give a total amplitude

$$m_{\text{tot}}^2 = 2N_c \int \frac{d^4 k_4}{(2\pi)^4} (2\pi R y_t)^2 \frac{\coth[k_4 \pi R]}{4} (\tanh[k_4 \pi R] - \coth[k_4 \pi R]). \quad (23)$$

We can rewrite this as

$$m_{\text{tot}}^2 = -\frac{N_c y_t^2}{4R^2} \int_0^\infty dx \frac{x^3}{\sinh^2[\pi x]} = -\frac{3\zeta(3)}{8\pi^4} \frac{N_c y_t^2}{R^2}, \quad (24)$$

precisely what we found before.

The finiteness is now easily understood: since loops originate on the Yukawa brane, any loops in the vicinity of the Yukawa brane are supersymmetric and cancel. Loops which probe the bulk, and reflect off the supersymmetry breaking brane will not cancel. However, these must vanish in the UV, leaving a finite result. We see this schematically in figure 2.

In fact, with this understanding, we need not have calculated the fermionic loop at all! One can simply calculate the scalar/F-component diagram, and subtract off the value in the $m \rightarrow 0$ limit, which yields eq. 24 again. This verifies our understanding of the finite result.

4.1 Scherk-Schwarz Theories

We have seen that supersymmetry breaking, localized away from the Yukawa brane can lead to UV-finite contributions. However, another alternative exists, utilizing the *global* properties of the five-dimensional space, via the Scherk-Schwarz mechanism [14].

Models can be constructed which give different boundary conditions to standard model particles and their superpartners (for instance, [1, 11, 15]). We will restrict our brief discussion to the model of [1] in which R -parity assignments classify which modes are even and which are odd.

Under R parity, various superfields transform as

$$\begin{aligned} X(x, y, \theta) &\rightarrow -X(x, y, -\theta), & X^c(x, y, \theta) &\rightarrow -X^c(x, y, -\theta), \\ H(x, y, \theta) &\rightarrow H(x, y, -\theta), & H^c(x, y, \theta) &\rightarrow H^c(x, y, -\theta), \\ V(x, y, \theta) &\rightarrow V(x, y, -\theta), & \Sigma(x, y, \theta) &\rightarrow \Sigma(x, y, -\theta), \end{aligned}$$

where X and H represent Q, U, D, L, E and H_u, H_d , respectively. It should be understood that the Higgs field need not propagate in the bulk, but we have listed its transformation properties in the event that it does.

Fermionic tops then have the ordinary Kaluza-Klein tower, $m = 0, R^{-1}, 2R^{-1}$, etc., and fermionic top conjugates remain unchanged from the initial orbifold, $m = R^{-1}, 2R^{-1}, 3R^{-1}$, etc. In contrast, stops and conjugate stops are shifted one unit up and down, respectively, $m = R^{-1}/2, 3R^{-1}/2, 5R^{-1}/2$, etc.

We can again calculate the contribution to the Higgs soft mass, and summing the entire tower yields the finite result

$$m_{\phi_{Hu}}^2 = -\frac{21\zeta(3)}{32\pi^4} N_c y_t^2 M_c^2. \quad (25)$$

4.2 Five dimensional interpretation

Again, we have summed the entire tower of states. Is this the correct thing to do? In the previous example, we could understand the finite behavior because of the locality of the supersymmetry breaking. What is the equivalent five dimensional understanding here?

The key point is that a Scherk-Schwarz breaking of supersymmetry is a *global* feature of the space, not a local one. Thus, short distance (i.e., high four-momentum) loops will only probe the local features of the spacetime, meaning those short distance loops should retain supersymmetric cancellations.

In mixed position-momentum space, this arises because of the different properties of particles and their superpartners as we continue around the extra dimension. That is, under $y \rightarrow y + 2\pi R$, the squarks will pick up a minus sign in their propagator.

$$\tilde{G}_\phi(k_4, y) = \sum_{n=-\infty}^{\infty} \frac{1}{2k_4} (-1)^n e^{-k_4|y-2\pi nR|}. \quad (26)$$

At the same time, the fermion picks up a plus sign when $y \rightarrow y + 2\pi R$.

$$\tilde{G}_\psi(k_4, y) = \sum_{n=-\infty}^{\infty} \frac{k_4}{2k_4} e^{-k_4|y-2\pi nR|}. \quad (27)$$

Note that the supersymmetry breaking occurs in these propagators only when the scalar and fermion both propagate *completely around the extra dimension* (i.e., $n \neq 0$). High four-momentum loops will exponentially damp out in this region, leaving only the supersymmetrically cancelling pieces of the propagators. A full calculation has been given in [1]. From this formulation, it is clear that the final result must be finite in the UV.

5 Phenomenology

We shall not delve deeply into the phenomenology of these models, but only comment briefly on their properties. Most notably, the squarks and sleptons are all degenerate up to corrections from

electroweak symmetry breaking. The gauginos are degenerate with the sfermions in the Scherk-Schwarz case, and can be degenerate, but are not necessarily so, in the localized supersymmetry breaking case.

Perhaps the most exciting change in phenomenology from standard SUSY theories is the change in the Higgs sector. If the Higgsino is the NLSP (with localized supersymmetry breaking, the gravitino is the LSP), and if the gluino is light enough to be produced at hadron colliders, it will decay $\tilde{g} \rightarrow \bar{q}q\tilde{h}$, followed by $\tilde{h} \rightarrow h\tilde{G}$. If the gluino is produced in large enough quantities, this could be the dominant mechanism for Higgs production.

So far we have not given a mechanism for generating μ . There are many possibilities. An attractive example is to utilize a next to minimal sector as in the NMSSM with a superpotential

$$W_{\text{Higgs}} = \lambda_H S H_u H_d + \lambda_B S B \bar{B} + \frac{\lambda_S}{3} S^3. \quad (28)$$

where S is a brane field and B, \bar{B} are bulk fields. Just as the Higgs soft mass is driven negative, so is the S soft mass, giving it a vev and generating a μ term. Ordinarily, one must be concerned that the trilinear couplings will run strong before the Planck scale. Here, the cutoff is only a few TeV, so we have great freedom in choosing the λ 's. Because of this, it is quite easy to have a Higgs much heavier than the $O(130\text{GeV})$ supersymmetric Higgs.

Precision electroweak studies have shown that a heavy Higgs is disfavored. However, in [24] it was shown that the inclusion of two nonrenormalizable operators, $(H^\dagger \tau^a H) W_{\mu\nu}^a B_{\mu\nu}$ and $|H^\dagger D_\mu H|^2$, could “fake” the existence of a light Higgs boson, at least with regard to precision electroweak studies. In [25] it was pointed out that if one allowed all possible operators uncontrolled by flavor symmetries at the same level (specifically $(H^\dagger D_\mu \tau^a H)(\bar{X} \gamma_\mu \tau^a X)_{X \rightarrow L, Q}$, $(H^\dagger D_\mu H)(\bar{X} \gamma_\mu X)_{X \rightarrow Q, U, D, LE}$, and $(\bar{L} \gamma_\mu \tau^a L)^2$), then to fake a light Higgs was incompatible with other precision studies. Put another way: any model that seeks to explain the precision studies of m_H with these nonrenormalizable operators must explain why the operators of [24] are anomalously large when compared with a host of other operators.

However, we note that all the operators which need be suppressed involve fields which live in the bulk. Thus, these operators are all naturally suppressed by a volume factor $1/2\pi R$ relative to the operators discussed by Hall and Kolda. The operator involving only the gauge bosons and the Higgs does not receive volume suppression, which simply goes into defining the four-dimensional gauge couplings from the five-dimensional gauge couplings. Thus, here we have precisely an example of a model in which the “bad” operators are automatically suppressed, while the “good” operators remain large. It is then quite natural in models where matter fields propagate in the bulk but the Higgs lies on a brane to give the appearance in precision tests that the Higgs is lighter than it is.

6 Conclusions

In theories in which the cutoff scale is $O(\text{TeV})$, we cannot truly say we have understood the hierarchy problem without understanding the origin of electroweak symmetry breaking. Here we have noted that with additional dimension of size $R^{-1} \sim \text{TeV}$ and five-dimensional supersymmetry, there can be a natural understanding of electroweak symmetry breaking. In this mechanism, the Higgs mass is calculably finite, and a loop factor below the masses of the superpartners, ameliorating fine-tuning issues.

The loops contributing to the Higgs soft mass have an extreme UV softness. The five-dimensional understanding of these models gives us an intuitive understanding - and calculationally tractable method - of understanding the origin of this behavior.

The phenomenology of these models is rich, containing heavy, degenerate superpartners. The Higgs sector can be radically changed. Not only are the theoretical controls on the light Higgs boson mass lifted, so, too, are those constraints from precision electroweak measurements. This makes the direct search for the Higgs boson in regions $M_H > 140\text{GeV}$ especially interesting.

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